

Baryon properties from light-front holographic QCD

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We investigate the properties of octet and decuplet baryons in a light-front holographic model. By taking into account the effect of nonvanishing quark mass, we obtain the modified light-front wave functions which are applicable at both low and high energy scales. We calculate the spectra, form factors, magnetic moments and electromagnetic radii of octet and decuplet baryons with the results all matching the experiments well. The axial charge, which describes the contribution of quark helicity to the proton spin in the quark-parton model at the high energy scale, is also consistent with the experimental value. Therefore, the light-front holographic method is successful in studying hadronic physics at all energy scales, and the nonzero quark mass is essential to understand the spin structures together with other low energy properties.

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Quantum chromodynamics (QCD) provides a fundamental description of strong interaction in terms of quark and gluon degrees of freedom, and has been proven successful in the high energy region where the perturbative effect dominates. It is still challenging to directly calculate the hadron properties at the low energy scale where the strong and nonlinear coupling between quarks and gluons plays the essential role. The light-front holography provides a new method to handle the strong interaction from basic considerations and has led to many remarkable results [1]. The light-front holography is based on the correspondence [2–4] between string states defined on the five-dimensional anti-de Sitter (AdS) space-time and conformal field theories (CFT) in physical space-time. Although the conformal symmetry of the classical QCD Lagrangian with massless quarks is broken by quantum effects, it is nearly conformal at high energy or short distance because of its asymptotic freedom [5]. There is also evidence from lattice QCD [6], the Dyson-Schwinger equation [7], empirical effective charges [8] and theoretical arguments [9] that the strong coupling constant has an infrared fixed point. Therefore, the AdS/CFT correspondence can be used to obtain a first approximation to QCD. This kind of methods have been applied to the hadronic spectrum [10], hard scattering [11] and strongly coupled quark-gluon plasma [12].

During the last decade, a connection is established between the conformal quantum mechanics and the light-front dynamics [13, 14], which provides a natural framework to reconcile the quark-parton model with QCD [15, 16]. The simple vacuum in the light-front quantization allows an unambiguous definition of the constituents of hadrons. All hadronic properties and partonic structures are encoded in the frame independent light-front wave functions (LFWFs). Therefore, solving the LFWFs becomes a central as well as challenging issue

in hadronic physics. Recently, an exact correspondence is found between the fifth-dimensional coordinate z in AdS space and the weighted impact separation variable ζ in physical space-time [17, 18]. This endows a clear physical meaning to the holographic variable. Some LFWFs are derived from some holographic models as a first approximation [19, 20].

In this paper, we study the baryon properties with the light-front holographic approach. Including the effects from nonzero quark mass, which explicitly breaks the conformal symmetry but plays an important role in understanding spin-related issues, we obtain the LFWFs with different orbital angular momenta and find that the spectra, magnetic moments, form factors and electromagnetic radii of octet and decuplet baryons are all well described. It is remarkable that the axial charge which describes the fraction of the quark helicity contribution to the proton spin is also consistent with the experimental value, when the quark mass effect is taken into account. Therefore, the new LFWFs we obtained are applicable to study the baryon properties at both low energy and high energy scales.

Hadrons, as bound states of the strong interaction, are the eigenstates of the QCD light-front Hamiltonian $H_{\text{LF}} = 2P^+P^- - \mathbf{P}_\perp^2$ with mass squares as the eigenvalues. Quantized at fixed light-front time, it can be expanded on the Fock state basis as

$$|H\rangle = \sum_n \int [dx][d^2\mathbf{k}_\perp] \psi_{n/H}(x_i, \mathbf{k}_{i\perp}) |n : x_i, \mathbf{k}_{i\perp}, \lambda_i\rangle, \quad (1)$$

where the integral measures are defined as

$$\int [dx] = \prod_{i=1}^n \int dx_i \delta(1 - \sum_{j=1}^n x_j), \quad (2)$$

$$\int [d^2\mathbf{k}_\perp] = \prod_{i=1}^n \int \frac{d^2\mathbf{k}_{i\perp}}{2(2\pi)^3} 16\pi^3 \delta^{(2)}(\sum_{j=1}^n \mathbf{k}_{j\perp}), \quad (3)$$

and λ_i represents the helicity and other internal degrees

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of freedom. Since there is an explicit separation of kinematical and dynamical terms in QCD light-front Hamiltonian, we can express the mass square of a hadron in terms of the LFWFs as

$$M_H^2 = \int [dx][d^2\mathbf{k}_\perp] \sum_{i=1}^n \frac{\mathbf{k}_{i\perp}^2 + m_i^2}{x_i} |\psi(x_i, \mathbf{k}_{i\perp})|^2 + \int [dx][d^2\mathbf{k}_\perp] \psi^*(x_i, \mathbf{k}_{i\perp}) U \psi(x_i, \mathbf{k}_{i\perp}), \quad (4)$$

where U is an effective potential. In this semiclassical approximation, some nondiagonal effects are neglected.

Using the Fourier transformation, one may obtain the expression in the coordinate space. For a two-body system, which means the hadron is regarded as an active quark and a spectator cluster, the eigenequation with massless constituents is expressed as [20]

$$\left(-\frac{d^2}{d\zeta^2} - \frac{1-4L^2}{4\zeta^2} + \tilde{U} \right) \phi(\zeta) = M_H^2 \phi(\zeta), \quad (5)$$

where the light-front variable $\zeta = \sqrt{x(1-x)}|\mathbf{b}_\perp|$ measures the separation between the quark and the spectator. It corresponds to the holographic variable z in AdS space. The transverse mode $\phi(\zeta)$ of the LFWF is defined as [17]

$$\psi(x, \mathbf{b}_\perp) = e^{iL\varphi} X(x) \frac{\phi(\zeta)}{\sqrt{2\pi\zeta}}. \quad (6)$$

The effective potential \tilde{U} is usually derived from the deformation of the AdS space, i.e., the soft-wall model [10], by introducing a dilaton $\varphi(z) = \lambda z^2$. For fermions, however, the dilaton can always be absorbed through the rescaling of the fields [21]. Therefore, an effective interaction $\rho(z)$ is introduced to the effective action in AdS space [22],

$$S_{\text{eff}} = \int d^4x dz \sqrt{g} e^{\varphi(z)} \left[\bar{\Psi} (i\Gamma^A e_A^M D_{M-\mu-\rho(z)}) \Psi + h.c. \right], \quad (7)$$

where Ψ is the bulk field, Γ^A is the tangent-space Dirac matrices, and e_A^M is the inverse vielbein.

The bayron light-front wave function satisfies the coupled linear differential equations [23],

$$\left(-\frac{d}{d\zeta} - \frac{\nu + \frac{1}{2}}{\zeta} - V(\zeta) \right) \phi_- = M \phi_+, \quad (8)$$

$$\left(\frac{d}{d\zeta} - \frac{\nu + \frac{1}{2}}{\zeta} - V(\zeta) \right) \phi_+ = M \phi_-, \quad (9)$$

where the subscripts \pm represent the chiral components. The confinement potential $V(\zeta)$ is determined by the effective interaction $\rho(z)$ as

$$V(\zeta) = \frac{R}{\zeta} \rho(\zeta). \quad (10)$$

In this study, we choose a linear potential $V(\zeta) = \lambda\zeta$ which can reproduce the Regge behavior for bayrons [23, 24]. This form can also be uniquely determined in the framework of superconformal algebra [25]. From the coupled differential equations (8) and (9), one can obtain the equivalent second-order equations as

$$\left(-\frac{d^2}{d\zeta^2} - \frac{1-4\nu^2}{4\zeta^2} + \tilde{U}_+(\zeta) \right) \phi_+ = M^2 \phi_+, \quad (11)$$

$$\left(-\frac{d^2}{d\zeta^2} - \frac{1-4(\nu+1)^2}{4\zeta^2} + \tilde{U}_-(\zeta) \right) \phi_- = M^2 \phi_-, \quad (12)$$

where the effective potential is expressed as

$$\tilde{U}_\pm(\zeta) = V^2(\zeta) \pm V'(\zeta) + \frac{1+2\nu}{\zeta} V(\zeta). \quad (13)$$

Comparing with (5), one may relate ν to the orbital angular momentum L . To have the separation between the kinematical and dynamical effects, we assign $L = \nu$ for the right-hand component and $L = \nu+1$ for the left-hand component.

We take the lowest energy solutions with $\nu = 0$ for the ground state bayrons,

$$\phi_+(\zeta) \sim \sqrt{2\zeta} e^{-\frac{\lambda\zeta^2}{2}}, \quad \phi_-(\zeta) \sim 2\zeta^{\frac{3}{2}} e^{-\frac{\lambda\zeta^2}{2}}. \quad (14)$$

To include the correction from nonzero masses of the constituents, we adopt the replacement [1],

$$\frac{\mathbf{k}_\perp^2}{x(1-x)} \rightarrow \frac{\mathbf{k}_\perp^2 + m_1^2}{x} + \frac{\mathbf{k}_\perp^2 + m_2^2}{1-x}, \quad (15)$$

to the LFWFs in the momentum space, where m_1 and m_2 are the effective masses of the quark and the spectator cluster respectively. Then the LFWFs have the form

$$\psi_+(x, \mathbf{k}_\perp) \sim \frac{4\pi}{\sqrt{\lambda x(1-x)}} e^{-\frac{1}{2\lambda} \left(\frac{\mathbf{k}_\perp^2 + m_1^2}{x} + \frac{\mathbf{k}_\perp^2 + m_2^2}{1-x} \right)}, \quad (16)$$

$$\psi_-(x, \mathbf{k}_\perp) \sim \frac{4\pi|\mathbf{k}_\perp|}{\lambda x(1-x)} e^{-\frac{1}{2\lambda} \left(\frac{\mathbf{k}_\perp^2 + m_1^2}{x} + \frac{\mathbf{k}_\perp^2 + m_2^2}{1-x} \right)}. \quad (17)$$

They actually describe $L = 0$ and $L = 1$ states respectively. In other words, the ground state bayrons contain the nonzero orbital angular momentum state in the light-front dynamics. This is a relativistic effect that can be physically understood from the Wigner rotation effect [26], which plays an important role in explaining the proton spin puzzle [27]. On the limit of zero quark mass, there is an equal probability to find the $L = 0$ and the $L = 1$ states in a baryon state [22]. In such a situation, the proton spin is all from the orbital angular momentum.

However, the nonvanishing quark mass changes the weights of these two states. Comparing the ratio with the generic ansatz,

$$\frac{|\psi_{L=0}|^2}{|\psi_{L=1}|^2} = \frac{(m_1 + x\mathcal{M})^2}{\mathbf{k}_\perp^2}, \quad (18)$$

where \mathcal{M} is a quantity with the dimension of mass, we identify the $x\mathcal{M}$ with $\sqrt{\lambda x(1-x)}$ in order to reproduce the ratio between (16) and (17) on the limit of massless quarks. We obtain the modified LFWFs as

$$\psi_{L=0}(x, \mathbf{k}_\perp) = N \frac{4\pi[m_1 + \sqrt{\lambda x(1-x)}]}{\lambda x(1-x)} \times e^{-\frac{1}{2\lambda}(\frac{\mathbf{k}_\perp^2 + m_1^2}{x} + \frac{\mathbf{k}_\perp^2 + m_2^2}{1-x})}, \quad (19)$$

$$\psi_{L=1}(x, \mathbf{k}_\perp) = N \frac{-4\pi(k^1 + ik^2)}{\lambda x(1-x)} \times e^{-\frac{1}{2\lambda}(\frac{\mathbf{k}_\perp^2 + m_1^2}{x} + \frac{\mathbf{k}_\perp^2 + m_2^2}{1-x})}, \quad (20)$$

where N is the normalization factor. Then the baryon mass can be evaluated via (4) with the contributions from kinematical energy, the confinement potential and constituent masses.

In the light-front formalism, the Dirac and Pauli form factors correspond to the spin-conserving and the spin-flip current matrix elements respectively [28]. With the electromagnetic current $V(z, Q^2)$ in AdS space, the Dirac form factor is expressed as

$$F_1(Q^2) = \sum_L \int dz \frac{R^4}{z^4} V(z, Q^2) \Psi_L^2(z). \quad (21)$$

Here we adopt the dressed current derived from the soft-wall model [17, 29],

$$V(z, Q^2) = \Gamma \left(1 + \frac{Q^2}{4\lambda}\right) U \left(\frac{Q^2}{4\lambda}, 0, \lambda z^2\right), \quad (22)$$

where $U(\alpha, \gamma, z)$ is the Tricomi confluent hypergeometric function.

Since the precise mapping of the Pauli form factor has not been carried out in holographic methods, a nonminimal electromagnetic coupling with the anomalous gauge invariant term is proposed [24],

$$\eta \int d^4x dz \sqrt{g} \bar{\Psi} e_M^A e_N^B [\Gamma_A, \Gamma_B] F^{MN} \Psi, \quad (23)$$

where η is an effective coupling constant. Then the Pauli form factor is expressed as

$$F_2(Q^2) = \eta \int dz \frac{R^3}{z^3} \bar{\Psi}_{L=0}(z) V(z, Q^2) \Psi_{L=1}(z). \quad (24)$$

The Sachs form factors are defined in terms of Dirac and Pauli form factors as [30]

$$G_E(Q^2) = F_1(Q^2) - \frac{Q^2}{4M^2} F_2(Q^2), \quad (25)$$

$$G_M(Q^2) = F_1(Q^2) + F_2(Q^2). \quad (26)$$

Fitted to the data, the coupling constant is chosen as $\eta = 1.3$, and the parameter in the effective potential is chosen as $\lambda_S = 0.105 \text{ GeV}^2$ for the scalar spectator

TABLE I. Effective masses of quarks and spectator clusters in the unit of MeV. The q represents the u quark and d quark.

m_q	50	$m_S(qq)$	130	$m_V(qq)$	770
m_s	203	$m_S(qs)$	554	$m_V(qs)$	1040
				$m_V(ss)$	1208

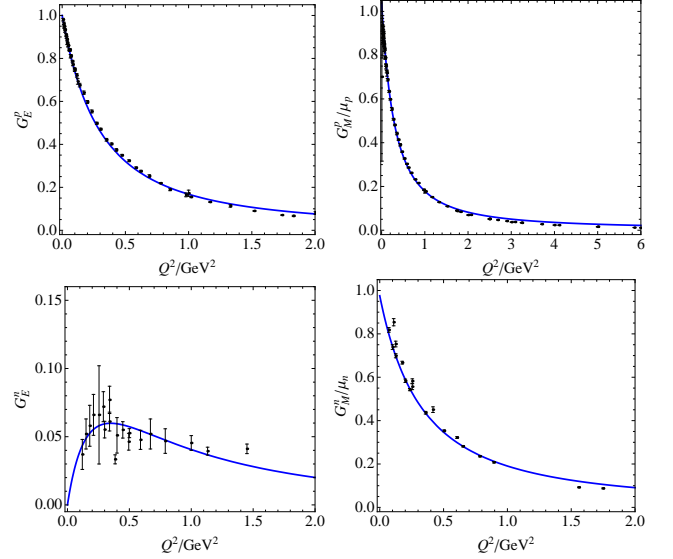


FIG. 1. The electromagnetic form factors of the proton and neutron. The curves are our model results. The data are taken from Refs. [31–41].

cluster and $\lambda_V = 0.170 \text{ GeV}^2$ for the vector spectator cluster. One may also use different η for the scalar and the vector spectator cases [24]. The mass parameters are listed in Table I. The results of the proton and the neutron form factors are plotted in Fig. 1 compared with the experimental data taken from Refs. [31–41].

We calculate the spectra of octet and decuplet baryons. The results are listed in Table II compared with the experimental values in Ref. [42]. The spectator masses $m_S(qq)$, $m_V(qq)$, $m_S(qs)$, $m_V(qs)$ and $m_V(ss)$ are close to those of the mesons π , ρ , K , K^* and ϕ . The N - Δ mass difference is obtained from the different parameters for the scalar and the vector spectators. An alternative treatment is to adopt different assignments to ν with L and $L + 1/2$ for N and Δ as used in Ref. [1].

The magnetic moment can be defined from the magnetic form factor on the limit of the zero momentum square transferred as $\mu = G_M(0)$. Our results of magnetic moments of octet baryons are listed in Table III compared with the experimental values in Ref. [42]. The radius is defined from the derivative of the form factor at zero momentum square transferred point as

$$\langle r^2 \rangle = -\frac{6}{F(0)} \frac{dF(Q^2)}{dQ^2} \Big|_{Q^2=0}. \quad (27)$$

The model results of the charge and the magnetic radii

TABLE II. The spectra of octet and decuplet baryons.

Baryon	M_{th}/MeV	M_{ex}/MeV [42]
N	977.8	p : 938.272046 ± 0.000021 n : 939.565379 ± 0.000021
Λ	1143	Λ^0 : 1115.683 ± 0.006
Σ	1146	Σ^+ : 1189.37 ± 0.07 Σ^0 : 1192.642 ± 0.024 Σ^- : 1197.449 ± 0.030
Ξ	1349	Ξ^0 : 1314.86 ± 0.20 Ξ^- : 1321.71 ± 0.07
Δ	1201	Δ : $1209 \sim 1210$
Σ^*	1370	Σ^{*+} : 1382.80 ± 0.35 Σ^{*0} : 1383.7 ± 1.0 Σ^{*-} : 1387.2 ± 0.5
Ξ^*	1525	Ξ^{*0} : 1531.80 ± 0.32 Ξ^{*-} : 1535.0 ± 0.6
Ω	1660	Ω^- : 1672.45 ± 0.29

TABLE III. The magnetic moments of octet baryons. The nuclear magneton is defined as $\mu_N = e\hbar/2m_p$.

Baryon	μ_{th}/μ_N	μ_{ex}/μ_N [42]
p	2.785	$2.792847356 \pm 0.000000023$
n	-1.790	-1.9130427 ± 0.0000005
Λ	-0.799	-0.613 ± 0.004
Σ^+	2.414	2.458 ± 0.010
Σ^-	-0.896	-1.160 ± 0.025
$\Sigma^0 \rightarrow \Lambda$	1.594	1.61 ± 0.08
Ξ^0	-1.181	-1.250 ± 0.014
Ξ^-	-0.589	-0.6507 ± 0.0025

of the nucleons are listed in Table IV compared with the experimental values in Ref. [42].

We need to emphasize that we calculate the flavor singlet axial charge which is interpreted as the contribution of quark helicities to the proton spin in the partonic language. Known as the “proton spin crisis,” it is often difficult to understand the helicity contribution at the high energy scale with other low energy baryon properties. In this study, by taking into account the quark mass effect, we obtain the value of the axial charge as 0.308 , which is consistent with the recent experimental analysis $0.330 \pm 0.011(\text{theo.}) \pm 0.025(\text{exp.}) \pm 0.028(\text{evol.})$ [43]. If the quark mass is neglected, the value will reduce to zero, which is an ultrarelativistic limit. Therefore, the

nonzero quark mass is important to describe hadronic spin structures.

TABLE IV. The charge and the magnetic radii of the proton and the neutron.

	Theory value	Experiment value [42]
r_E^p	0.8819 fm	0.8775 ± 0.0051 fm
r_M^p	0.8783 fm	0.777 ± 0.016 fm
$\langle (r_E^n)^2 \rangle$	-0.1390 fm^2	$-0.1161 \pm 0.0022 \text{ fm}^2$
r_M^n	0.8392 fm	$0.862^{+0.009}_{-0.008}$ fm

In summary, we investigate the octet and decuplet baryon properties in a light-front holographic soft-wall model. On the basis of the correspondence established on the limit of massless quarks, we include the corrections from nonzero masses of the constituents and obtain the modified LFWFs of both $L = 0$ and $L = 1$ states for baryons. Our results of baryon spectra, form factors, magnetic moments, and the charge and the magnetic radii match the experimental data well. The axial charge, which describes the quark helicity contribution to the nucleon spin, is also consistent with the analysis from the experimental data. Therefore, the AdS/CFT correspondence, or more generally the gauge/gravity duality, provides a powerful tool to study hadronic physics at both low and high energy scales, and if the effect from nonvanishing quark mass is taken into account, it is also successful in the study of spin-related issues.

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Note added.— Some phenomenological baryon LFWFs based on the soft-wall model with finite quark mass were also proposed in Ref. [44] and were applied to investigate nucleon structures.

- [1] S. J. Brodsky, G. F. de Teramond, H. G. Dosch, and J. Erlich, Phys. Rep. **584**, 1 (2015).
- [2] J. M. Maldacena, Int. J. Theor. Phys. **38**, 1113 (1999) [Adv. Theor. Math. Phys. **2**, 231 (1998)].
- [3] S. S. Gubser, I. R. Klebanov, and A. M. Polyakov, Phys. Lett. B **428**, 105 (1998).
- [4] E. Witten, Adv. Theor. Math. Phys. **2**, 253 (1998).
- [5] D. J. Gross and F. Wilczek, Phys. Rev. D **8**, 3633 (1973); D. J. Gross and F. Wilczek, Phys. Rev. Lett. **30**, 1343

- (1973); H. D. Politzer, Phys. Rev. Lett. **30**, 1346 (1973).
- [6] S. Furui and H. Nakajima, Phys. Rev. D **76**, 054509 (2007).
- [7] L. von Smekal, A. Hauck, and R. Alkofer, Phys. Rev. Lett. **79**, 3591 (1997); D. Binosi and J. Papavassiliou, Phys. Rep. **479**, 1 (2009).
- [8] A. Deur, V. Burkert, J.-P. Chen, and W. Korsch, Phys. Lett. B **650**, 244 (2007); A. Deur, V. Burkert, J.-P. Chen, and W. Korsch, Phys. Lett. B **665**, 349 (2008).

- [9] S. J. Brodsky and R. Shrock, Phys. Lett. B **666**, 95 (2008).
- [10] A. Karch, E. Katz, D. T. Son, and M. A. Stephanov, Phys. Rev. D **74**, 015005 (2006).
- [11] J. Polchinski and M. J. Strassler, Phys. Rev. Lett. **88**, 031601 (2002).
- [12] J. Casalderrey-Solana, H. Liu, D. Mateos, K. Rajagopal, and U. A. Wiedemann, arXiv:1101.0618.
- [13] S. J. Brodsky and G. F. de Teramond, Phys. Lett. B **582**, 211 (2004).
- [14] S. J. Brodsky, G. F. De Teramond, and H. G. Dosch, Phys. Lett. B **729**, 3 (2014).
- [15] S. J. Brodsky, H. C. Pauli, and S. S. Pinsky, Phys. Rep. **301**, 299 (1998).
- [16] P. A. M. Dirac, Rev. Mod. Phys. **21**, 392 (1949).
- [17] S. J. Brodsky and G. F. de Teramond, Phys. Rev. D **77**, 056007 (2008).
- [18] S. J. Brodsky and G. F. de Teramond, Phys. Rev. D **78**, 025032 (2008).
- [19] S. J. Brodsky and G. F. de Teramond, Phys. Rev. Lett. **96**, 201601 (2006).
- [20] G. F. de Teramond and S. J. Brodsky, Phys. Rev. Lett. **102**, 081601 (2009).
- [21] I. Kirsch, J. High Energy Phys. **09** (2006) 052.
- [22] G. F. de Teramond, H. G. Dosch, and S. J. Brodsky, Phys. Rev. D **87**, 075005 (2013).
- [23] S. J. Brodsky and G. F. de Teramond, Proceedings of the International School of Subnuclear Physics, Vol. **45**, 139 (2009), arXiv:0802.0514.
- [24] Z. Abidin and C. E. Carlson, Phys. Rev. D **79**, 115003 (2009).
- [25] G. F. de Teramond, H. G. Dosch, and S. J. Brodsky, Phys. Rev. D **91**, 045040 (2015); H. G. Dosch, G. F. de Teramond, and S. J. Brodsky, Phys. Rev. D **91**, 085016 (2015).
- [26] E. P. Wigner, Annals Math. **40**, 149 (1939) [Nucl. Phys. Proc. Suppl. **6**, 9 (1989)].
- [27] B.-Q. Ma, J. Phys. G **17**, L53 (1991); B.-Q. Ma and Q.-R. Zhang, Z. Phys. C **58**, 479 (1993).
- [28] S. J. Brodsky and S. D. Drell, Phys. Rev. D **22**, 2236 (1980).
- [29] H. R. Grigoryan and A. V. Radyushkin, Phys. Rev. D **76**, 095007 (2007).
- [30] R. G. Sachs, Phys. Rev. **126**, 2256 (1962).
- [31] J. Arrington, W. Melnitchouk, and J. A. Tjon, Phys. Rev. C **76**, 035205 (2007).
- [32] R. C. Walker *et al.*, Phys. Rev. D **49**, 5671 (1994).
- [33] F. Borkowski, G. G. Simon, V. H. Walther, and R. D. Wendling, Nucl. Phys. **B93**, 461 (1975).
- [34] G. Kubon *et al.*, Phys. Lett. B **524**, 26 (2002).
- [35] D. I. Glazier *et al.*, Eur. Phys. J. A **24**, 101 (2005).
- [36] P. Markowitz *et al.*, Phys. Rev. C **48**, R5 (1993).
- [37] E. E. W. Bruins *et al.*, Phys. Rev. Lett. **75**, 21 (1995).
- [38] A. Lung *et al.*, Phys. Rev. Lett. **70**, 718 (1993).
- [39] H. Zhu *et al.* (E93026 Collaboration), Phys. Rev. Lett. **87**, 081801 (2001).
- [40] G. Warren *et al.* (Jefferson Lab E93-026 Collaboration), Phys. Rev. Lett. **92**, 042301 (2004).
- [41] D. Rohe *et al.*, Phys. Rev. Lett. **83**, 4257 (1999).
- [42] K. A. Olive *et al.* (Particle Data Group), Chin. Phys. C **38**, 090001 (2014).
- [43] A. Airapetian *et al.* (HERMES Collaboration), Phys. Rev. D **75**, 012007 (2007).
- [44] T. Gutsche, V. E. Lyubovitskij, I. Schmidt, and A. Vega, Phys. Rev. D **89**, 054033 (2014); T. Gutsche, V. E. Lyubovitskij, I. Schmidt, and A. Vega, Phys. Rev. D **91**, 054028 (2015).